DEGENERATE BESS

AT FUTURE e^+e^- COLLIDERS[*]

DANIELE DOMINICI

Dip. di Fisica, Univ. di Firenze, Sezione I.N.F.N. Firenze
Lgo E. Fermi 2, 50125 Firenze, Italy
E-mail: dominici@fi.infn.it

ABSTRACT

An effective lagrangian describing a strong interacting electroweak sector is considered. It contains new vector and axial-vector resonances all degenerate in mass and mixed with W and Z. The model, for large mass of these degenerate gauge bosons, becomes identical to the standard model in the limit of infinite Higgs mass. The limits on the parameter space of this model from future e^+e^- colliders are presented.

1. Introduction

The standard $SU(2) \otimes U(1)$ gauge theory of the electroweak interactions is in good agreement with the current experimental data, apart the 2 - 3 σ discrepancies in R_b and A_{LR}. Nevertheless there is no yet evidence for the mechanism which is responsible for the breakdown of the symmetry to the $U(1)$ electromagnetic. It is usually assumed that the breaking of the electroweak symmetry is due to the vacuum expectation value of some elementary scalar.

In this talk I would like to discuss a different option, a dynamical breaking of the electroweak symmetry: some new interaction induces a breaking at a scale Λ of order 1 TeV. Effective theories can be built on the basis of the low energy symmetry properties. We can build the low energy theory describing goldstones using the classical technique of Callan, Coleman, Wess and Zumino (CCWZ) [1], treating the pseudoscalars as the goldstone bosons of a spontaneously broken symmetry G to a subgroup H. In the simplest example a chiral symmetry $G = SU(2)_L \otimes SU(2)_R$ is broken to the diagonal subgroup $H = SU(2)_{L+R}$, producing three Goldstone bosons, which become via the Higgs mechanism the longitudinal degrees of freedom of W and Z. In general such a theory can contain also new resonances, like the ρ in QCD.

To build the effective low energy theory describing Goldstones and vectors, one can use the CCWZ non linear representations of a chiral symmetry G and consider-

---

ing (à la Weinberg [2]) the ρ as the gauge field of the unbroken symmetry group H. This theory is not renormalizable in the standard sense. We can order the terms in the lagrangian in an energy expansion according the number of derivatives and truncate at some finite order. The higher order terms will be proportional to the inverse power of the parameter Λ.

In a completely equivalent way one can use the *hidden gauge symmetry* approach. Theories with non linearly realized symmetry $G \to H$ can be linearly realized by enlarging the gauge symmetry G to $G \otimes H' \to H_D = diag(H \otimes H')$. H' is a local gauge group and the ρ is the gauge field associated to H' [3,4].

The BESS (Breaking Electroweak Symmetry Strongly) model was built in this way, using $G = SU(2)_L \otimes SU(2)_R$, $H = SU(2)_V$ [5], and considering the gauging of the $SU(2)_W \otimes U(1)_Y$. This model is an effective lagrangian parametrization of the electroweak symmetry breaking. A new triplet of vector bosons, mixed with W and Z, is present. The parameters of the BESS model are the mass M_V of these new bosons, their self coupling g'' and a third parameter b whose strength characterizes the direct coupling of V to the fermions. The new charged vector bosons can be studied in the channel $W^\pm Z \to l^\pm \nu 2l$ at LHC, after their Drell-Yan production from to the initial quarks, up to masses of the order of 2 TeV [6].

At the previous LCWs in Saariselkä [7] and in Waikoloa [8] I discussed how future $e^+ e^-$ colliders could restrict the parameter space of the BESS model.

In principle such a theory can also include axial vector resonances [9], like the a_1 of QCD, or can have a larger symmetry like $SU(8) \otimes SU(8)$ [10].

In this talk I will present the results of a new phenomenological analysis on a particular version of these models, based on a chiral $SU(2)_L \otimes SU(2)_R$, containing new vector and axial-vector particles degenerate in mass (degenerate BESS) [11]. This particular choice of the parameters corresponds to an enlarged symmetry, and implies that leading contribution (in the large M_V expansion) to the $\epsilon_{1,2,3}$ (or S, T, U) parameters is zero; therefore the model is not much constrained by the existing data. In degenerate BESS relatively light resonances are compatible with the electroweak data, as given by LEP and Tevatron.

2. The model

Let me firstly recall how the most general lagrangian up two derivatives for the linearly realized $SU(2)_L \otimes SU(2)_R \otimes SU(2)_V \to SU(2)$ symmetry is built. The coordinates of the manifold $G/H = SU(2)_L \otimes SU(2)_R/SU(2)$ are substituted by a group element $g = (L, R) \in G$. The Goldstones bosons are represented by two unitary matrices L and R whose transformations are

$$L \to g_L L h \quad R \to g_R R h$$

with $g_{L,R} \in SU(2)_{L,R}$ and $h \in SU(2)_V$. Using these fields one can reconstruct the field $U = LR^\dagger$ which transforms as $U \to g_L U g_R^\dagger$ and describes the usual field of $SU(2)_L \otimes SU(2)_R / SU(2)$. We introduce also a gauge field $V_\mu = \frac{i}{2} g'' \frac{\tau_i}{2} V_\mu^i$ in $LieSU(2)_V$ and build the covariant derivatives $D_\mu L = \partial_\mu L - L V_\mu$, $D_\mu R = \partial_\mu R - R V_\mu$. The leading terms in the effective lagrangian invariant with respect to $SU(2)_L \otimes SU(2)_R$ and $L \leftrightarrow R$ transformation are given by

$$L_{eff} = -\frac{v^2}{4} \left[Tr(L^\dagger D_\mu L - R^\dagger D_\mu R)^2 + \alpha Tr(L^\dagger D_\mu L + R^\dagger D_\mu R)^2 \right] + ...$$

where v and α are arbitrary parameter. Going into the unitary gauge $L = R^\dagger = \exp[i\tau_i \pi_i / (2v)]$ one gets an effective lagrangian describing goldstones and massive vector mesons with $M_V = v / 2g'' \sqrt{\alpha}$.

After the $SU(2)_W \otimes U(1)_Y$ gauging and the identification $v^2 = 1/(\sqrt{2} G_F)$, G_F being the Fermi constant, one get the BESS model [5].

This procedure can be extended to include also axial vector resonances. Let $G = SU(2)_L \otimes SU(2)_R$ and $H' = SU(2)_L \otimes SU(2)_R$. The nine Goldstone bosons resulting from the spontaneous breaking of $G' = G \otimes H'$ to H_D, can be described by three independent $SU(2)$ elements: L, R and M, with the following transformations properties

$$L' = g_L L h_L, \quad R' = g_R R h_R, \quad M' = h_R^\dagger M h_L \tag{2.1}$$

with $g_{L,R} \in SU(2)_{L,R} \subset G$ and $h_{L,R} \in H'$. Moreover we shall require the invariance under the discrete left-right transformation $P: \quad L \leftrightarrow R, \quad M \leftrightarrow M^\dagger$ which combined with the usual space inversion allows to build the parity transformation on the fields. If we ignore the transformations of eq.(2.1), the largest possible global symmetry of the low-energy theory is given by the requirement of maintaining for the transformed variables L', R' and M' the character of $SU(2)$ elements, or $G_{max} = [SU(2) \otimes SU(2)]^3$, consisting of three independent $SU(2) \otimes SU(2)$ factors, acting on each of the three variables separately. As we shall see, it happens that, for specific choices of the parameters of the theory, the symmetry G' gets enlarged to G_{max}.

The most general $G' \otimes P$ invariant lagrangian is given by [9]

$$L_G = -\frac{v^2}{4}[a_1 I_1 + a_2 I_2 + a_3 I_3 + a_4 I_4] \tag{2.2}$$

plus the kinetic terms L_{kin}. The four invariant terms I_i $(i = 1, ...4)$ are given by:

$$I_1 = tr[(V_0 - V_1 - V_2)^2] \quad I_2 = tr[(V_0 + V_2)^2] \quad I_3 = tr[(V_0 - V_2)^2] \quad I_4 = tr[V_1^2]$$

where

$$V_0^\mu = L^\dagger D^\mu L \quad V_1^\mu = M^\dagger D^\mu M \quad V_2^\mu = M^\dagger (R^\dagger D^\mu R) M$$

and the covariant derivatives are

$$D_\mu L = \partial_\mu L - L\tilde{L}_\mu \quad D_\mu R = \partial_\mu R - R\tilde{R}_\mu$$

$$D_\mu M = \partial_\mu M - M\tilde{L}_\mu + \tilde{R}_\mu M$$

where $\tilde{L}_\mu(\tilde{R}_\mu)$ are gauge fields of $SU(2)_{L(R)} \subset H'$ (instead of working with vector and axial-vector we work with these left and right combinations).

The kinetic terms are given by

$$L_{kin} = \frac{1}{g''^2} tr[F_{\mu\nu}(\tilde{L})]^2 + \frac{1}{g''^2} tr[F_{\mu\nu}(\tilde{R})]^2$$

where g'' is the gauge coupling constant for the gauge fields $\tilde{L}_\mu$ and $\tilde{R}_\mu$, and $F_{\mu\nu}(\tilde{L})$, $F_{\mu\nu}(\tilde{R})$ are the usual field tensors.

The model I will discuss is characterized by the following choice of parameters $a_4 = 0$, $a_2 = a_3$ [11]. In order to discuss the symmetry properties it is useful to observe that the invariant I_1 could be re-written as $I_1 = -tr(\partial_\mu U^\dagger \partial^\mu U)$ with $U = LM^\dagger R^\dagger$ and the lagrangian as

$$L_G = \frac{v^2}{4}\{a_1\ tr(\partial_\mu U^\dagger \partial^\mu U) + 2\ a_2\ [tr(D_\mu L^\dagger D^\mu L) + tr(D_\mu R^\dagger D^\mu R)]\} \qquad (2.3)$$

Each of the three terms in the above expressions is invariant under an independent $SU(2) \otimes SU(2)$ group

$$U' = \omega_L U \omega_R^\dagger, \quad L' = g_L L h_L, \quad R' = g_R R h_R$$

The overall symmetry is $G_{max} = [SU(2) \otimes SU(2)]^3$, with a part H' realized as a gauge symmetry. With the particular choice $a_4 = 0$, $a_3 = a_2$, as we see from eq.(2.3), the mixing between $\tilde{L}_\mu$ and $\tilde{R}_\mu$ is vanishing, and the new states are degenerate in mass. Moreover, as it follows from eq.(2.3), the longitudinal modes of the fields are entirely provided by the would-be Goldstone bosons in L and R. This means that the pseudoscalar particles remaining as physical states in the low-energy spectrum are those associated to U. They in turn can provide the longitudinal components to the W and Z particles, in an effective description of the electroweak breaking sector.

The peculiar feature of degenerate BESS is that the new bosons are not coupled to those Goldstone bosons which are absorbed to give mass to $W^\pm$ and Z. As a consequence the channels $W_L Z_L$ and $W_L W_L$ are not strongly enhanced as it usually happens in models with a strongly interacting symmetry breaking sector and this implies larger branching ratios of the new resonances into fermion pairs.

The coupling of the model to the electroweak $SU(2)_W \otimes U(1)_Y$ gauge fields is obtained via the minimal substitution

$$D_\mu L \to D_\mu L + \tilde{W}_\mu L \quad D_\mu R \to D_\mu R + \tilde{Y}_\mu R \quad D_\mu M \to D_\mu M$$

where

$$\tilde{W}_\mu = ig\tilde{W}_\mu^a \frac{\tau^a}{2} \quad \tilde{Y}_\mu = ig'\tilde{Y}_\mu \frac{\tau^3}{2}$$

$$\tilde{L}_\mu = i\frac{g''}{\sqrt{2}}\tilde{L}_\mu^a \frac{\tau^a}{2} \quad \tilde{R}_\mu = i\frac{g''}{\sqrt{2}}\tilde{R}_\mu^a \frac{\tau^a}{2}$$

with g, g' the $SU(2)_W \otimes U(1)_Y$ gauge coupling constant and τ^a the Pauli matrices.

By introducing the canonical kinetic terms for W_μ^a and Y_μ and going into the unitary gauge we get

$$\begin{aligned}
\mathcal{L} = &-\frac{v^2}{4}\left[a_1 tr(\tilde{W}_\mu - \tilde{Y}_\mu)^2 + 2a_2 tr(\tilde{W}_\mu - \tilde{L}_\mu)^2 + 2a_2 tr(\tilde{Y}_\mu - \tilde{R}_\mu)^2\right] \\
&+ \mathcal{L}^{kin}(\tilde{W}, \tilde{Y}, \tilde{L}, \tilde{R})
\end{aligned} \tag{2.4}$$

We have used tilded quantities to reserve untilded variables for mass eigenstates.

The standard model (SM) relations are obtained in the limit $g'' \gg g, g'$. Actually, for a very large g'', the kinetic terms for the fields $\tilde{L}_\mu$ and $\tilde{R}_\mu$ drop out, and $\mathcal{L}$ reduces to the first term in eq.(2.4). This term reproduces precisely the mass term for the ordinary gauge vector bosons in the SM, provided we assume $a_1 = 1$. Finally let us consider the fermions of the SM and denote them by ψ_L and ψ_R. They couple to $\tilde{L}$ and $\tilde{R}$ via the mixing with the standard $\tilde{W}$ and $\tilde{Y}$:

$$\begin{aligned}
\mathcal{L}_{fermion} = &\overline{\psi}_L i\gamma^\mu\left(\partial_\mu + ig\tilde{W}_\mu^a \frac{\tau^a}{2} + \frac{i}{2}g'(B-L)\tilde{Y}_\mu\right)\psi_L \\
&+ \overline{\psi}_R i\gamma^\mu\left(\partial_\mu + ig'\tilde{Y}_\mu \frac{\tau^3}{2} + \frac{i}{2}g'(B-L)\tilde{Y}_\mu\right)\psi_R
\end{aligned}$$

where $B(L)$ is the baryon (lepton) number, and $\psi = (\psi_u, \psi_d)$.

By separating the charged and the neutral gauge bosons the quadratic lagrangian is given by:

$$\begin{aligned}
\mathcal{L}^{(2)} = &\frac{v^2}{4}[(1+2a_2)g^2\tilde{W}_\mu^+\tilde{W}^{\mu-} + a_2 g''^2(\tilde{L}_\mu^+\tilde{L}^{\mu-} + \tilde{R}_\mu^+\tilde{R}^{\mu-}) \\
&- \sqrt{2}a_2 gg''(\tilde{W}_\mu^+\tilde{L}^{\mu-} + \tilde{W}_\mu^-\tilde{L}^{\mu+})] \\
&+ \frac{v^2}{8}[(1+2a_2)(g^2\tilde{W}_3^2 + g'^2\tilde{Y}^2) + a_2 g''^2(\tilde{L}_3^2 + \tilde{R}_3^2) \\
&- 2gg'\tilde{W}_{3\mu}\tilde{Y}^\mu - 2\sqrt{2}a_2 g''(g\tilde{W}_3\tilde{L}_3^\mu + g'\tilde{Y}_\mu\tilde{R}_3^\mu)]
\end{aligned} \tag{2.5}$$

Therefore the $R^{\pm}$ fields are unmixed and their mass can be easily read: $M_{R^{\pm}} \equiv M = vg''\sqrt{a_2}/2$. We will parametrize the model by using, in addition to the SM parameters, M and g/g''.
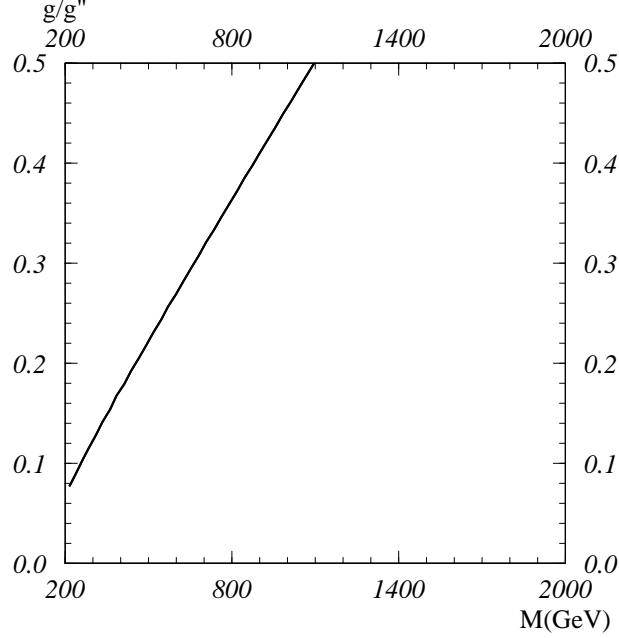

Fig. 1. 90% C.L. contour on the plane $(M, g/g'')$ obtained by comparing the values of the ϵ parameters from the model with the experimental data from LEP. The allowed region is below the curve.

Eigenvalues and eigenvectors for the remaining fields can be found in [11]. As already said the heavy fields have all the degenerate mass M in the large g'' limit. By using eq.(2.5) one can show that at the leading order in q^2/M^2 the contribution of the model to all ϵ parameters [12] is equal to zero [11]. This is due to the fact that in the $M \to \infty$ limit, this model decouples. We can perform the low-energy limit at the next-to-leading order and study the virtual effects of the heavy particles. Working at the first order in $1/g''^2$ we get $\epsilon_1 = -(c_\theta^4 + s_\theta^4)/(c_\theta^2)\ X$, $\epsilon_2 = -c_\theta^2\ X$, $\epsilon_3 = -X$ with $X = 2(M_Z^2/M^2)(g/g'')^2$. All these deviations are of order X which contains a double suppression factor M_Z^2/M^2 and $(g/g'')^2$. The sum of the SM contributions, functions of the top and Higgs masses, and of these deviations has to be compared with the experimental values for the ϵ parameters, determined from the all available LEP data and the M_W measurement at Tevatron [13]: $\epsilon_1 = (3.8 \pm 1.5) \cdot 10^{-3}$, $\epsilon_2 = (-6.4 \pm 4.2) \cdot 10^{-3}$, $\epsilon_3 = (4.6 \pm 1.5) \cdot 10^{-3}$. Taking into account the SM values $(\epsilon_1)_{SM} = 4.4 \cdot 10^{-3}$, $(\epsilon_2)_{SM} = -7.1 \cdot 10^{-3}$, $(\epsilon_3)_{SM} = 6.5 \cdot 10^{-3}$ for $m_{top} = 180\ GeV$ and $m_H = 1000\ GeV$, we find, from the combinations of the previous experimental results, the 90% C.L. limit on g/g'' versus the mass M given by the solid line in

Fig.1. The allowed region is the one below the continuous line.

3. Degenerate BESS at e^+e^- future colliders

In this section I will discuss the sensitivity of the model at LEP2 and future e^+e^- linear colliders, for different options of total centre of mass energies and luminosities.

Cross-sections and asymmetries for the channel $e^+e^- \rightarrow f^+f^-$ and $e^+e^- \rightarrow W^+W^-$ in the Standard Model and in the degenerate BESS model at tree level have been studied [11]. The BESS states relevant for the analysis at e^+e^- colliders are L_3 and R_3. Their coupling to fermions can be found in [11]. I will not consider the direct production of R_3 and L_3 from e^+e^-, but rather their indirect effects in the $e^+e^- \rightarrow f^+f^-$ and $e^+e^- \rightarrow W^+W^-$ cross-sections. In the fermion channel the study is based on the following observables: the total hadronic $(\mu^+\mu^-)$ cross-sections σ^h (σ^μ), the forward-backward and left-right asymmetries $A_{FB}^{e^+e^-\rightarrow\mu^+\mu^-}$, $A_{FB}^{e^+e^-\rightarrow\bar{b}b}$, $A_{LR}^{e^+e^-\rightarrow\mu^+\mu^-}$, $A_{LR}^{e^+e^-\rightarrow h}$ and $A_{LR}^{e^+e^-\rightarrow\bar{b}b}$. At LEP2 we can add to the previous observables the W mass measurement. The result of this analysis shows that LEP2 will not improve considerably the existing limits [14].
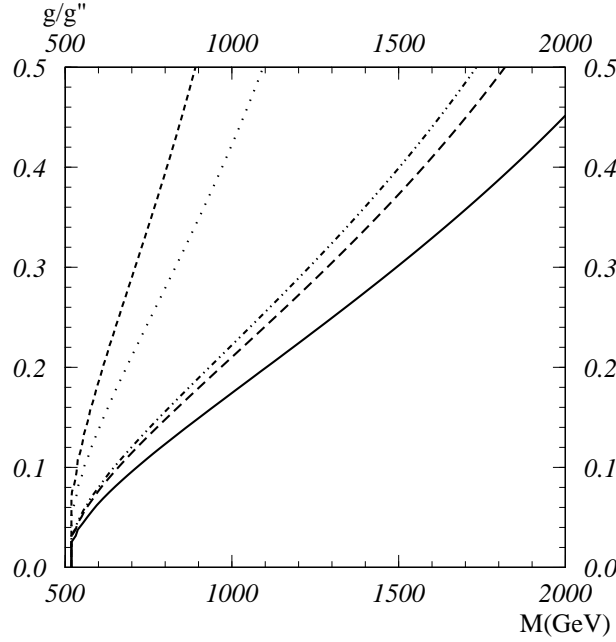

Fig. 2. 90% C.L. contour on the plane $(M, g/g'')$ from e^+e^- at $\sqrt{s} = 500 \ GeV$ with an integrated luminosity of $20 fb^{-1}$ from unpolarized observables. Allowed regions are below the curves. (Dashed-dotted σ^h, dashed σ^μ, dotted A_{FB}^μ, the uppermost dashed A_{FB}^b, continuous all combined).

To improve these limits it is necessary to consider higher energy colliders. Two options for a high energy e^+e^- collider have been studied: $\sqrt{s} = 500\ GeV$ ($\sqrt{s} = 1\ TeV$) with an integrated luminosity of $20fb^{-1}$ ($80fb^{-1}$).

In Fig. 2 we present the 90% C.L. contour on the plane $(M, g/g'')$ from e^+e^- at $\sqrt{s} = 500\ GeV$ with an integrated luminosity of $20fb^{-1}$ for various observables. The dashed-dotted line represents the limit from σ^h with an assumed relative error of 2%; the dashed line near to the preceeding one is σ^μ (relative error 1.3%), the dotted line is A^μ_{FB} (error 0.5%) and the uppermost dashed line is A^b_{FB} (error 0.9%).

As it is evident more stringent bounds come from the cross-section measurements. Asymmetries give less restrictive bounds due to a compensation between the L_3 and R_3 exchange. By combining all the deviations in the previously considered observables we get the limit shown by the continuous line.
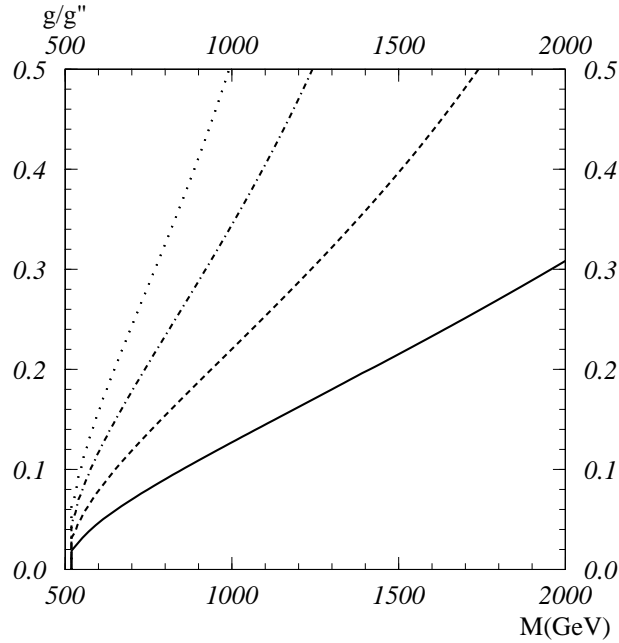

Fig. 3. 90% C.L. contour on the plane $(M, g/g'')$ from e^+e^- at $\sqrt{s} = 500\ GeV$ with an integrated luminosity of $20fb^{-1}$ from polarized observables. Allowed regions are below the curves. (Dashed-dotted A^μ_{LR}, dashed A^h_{LR}, dotted A^b_{LR}, continuous all unpolarized and polarized combined).

Polarized electron beams allow to get further limit in the parameter space as shown in Fig. 3. We neglect the error on the measurement of the polarization and use a polarization value equal to 0.5. The dashed-dotted line represents the limit from A^μ_{LR} (error 0.6%), the dashed line from A^h_{LR} (error 0.4%), the dotted line from A^b_{LR} (error 1.1%). Combining all the polarized and unpolarized beam observables we get the bound shown by the continuous line. In conclusion a substantial

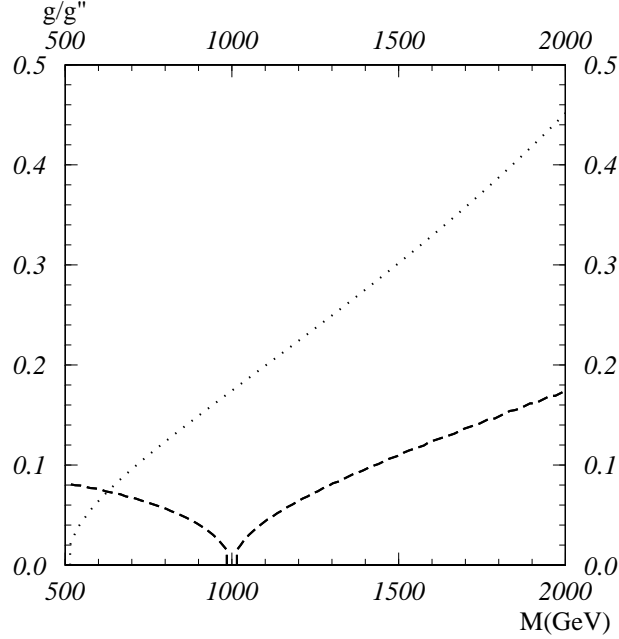improvement with respect to the LEP bounds, even without polarized beams is obtained.

Fig. 4. 90% C.L. contour on the plane $(M, g/g'')$ from e^+e^- at $\sqrt{s} = 500\ GeV$ with an integrated luminosity of $20fb^{-1}$ and $\sqrt{s} = 1000\ GeV$ with an integrated luminosity of $80fb^{-1}$. Allowed regions are below the curves.

In Fig. 4 a combined picture of the 90% C.L. contours on the plane $(M, g/g'')$ from e^+e^- at two values of $\sqrt{s}$ is shown. The dotted line represents the limit from the combined unpolarized observables at $\sqrt{s} = 500\ GeV$ with an integrated luminosity of $20fb^{-1}$; the dashed line is the limit from the combined unpolarized observables at $\sqrt{s} = 1000\ GeV$ with an integrated luminosity of $80fb^{-1}$. As expected increasing the energy of the collider and rescaling the integrated luminosity result in stronger bounds on the parameter space.

The WW final state, considering the observables given in [11] has been also studied. However the new channel does not modify the strong limits obtained using the fermion final state. This is because the degenerate model has no strong enhancement of the WW channel, present in the usual strong electroweak models.

4. Acknowledgements

this talk.

5. References

1. Coleman, J. Wess and B. Zumino, *Phys. Rev.* **177** (1969) 2239; C. G. Callan, S. Coleman, J. Wess and B. Zumino, *Phys. Rev.* **177** (1969) 2247.
2. S. Weinberg, *Phys. Rev.* **166** (1968) 1568.
3. A.P. Balachandran, A. Stern and G. Trahern, *Phys. Rev.* **D19** (1979) 2416.
4. M. Bando, T. Kugo, S. Uehara, K. Yamawaki and T. Yanagida, *Phys. Rev. Lett.* **54** (1985) 1215.
5. R. Casalbuoni, S. De Curtis, D. Dominici and R. Gatto, *Phys. Lett.* **B155** (1985) 95; and *Nucl. Phys.* **B282** (1987) 235.
6. R. Casalbuoni, P. Chiappetta, S. De Curtis, F. Feruglio, R. Gatto, B. Mele and J. Terron, *Phys. Lett.* **B249** (1991) 130. I. Josa, T. Rodrigo and F. Pauss, CERN 90-10, volume II, p. 796, Proceedings of Large Hadron Collider Workshop, Aachen, 4-9 October 1990, Eds. G. Jarlskog.
7. D.Dominici, in *Physics and Experiments with Linear Colliders*, eds. R. Orava, P. Eerola and M. Nordberg (World Scientific, Singapore, 1992), p. 509.
8. R. Casalbuoni, S. De Curtis, D. Dominici, P. Chiappetta, A. Deandrea and R. Gatto, in *Physics and Experiments with Linear e^+e^- Colliders*, eds. F. A. Harris, S. L. Olsen, S. Pakvasa and X. Tata (World Scientific, Singapore, 1993), p. 887.
9. R. Casalbuoni, S. De Curtis, D. Dominici F. Feruglio and R. Gatto, Int. Jour. Mod. Phys. **A4** (1989) 1065.
10. R. Casalbuoni, A. Deandrea, S. De Curtis, N. Di Bartolomeo, D. Dominici, F. Feruglio and R. Gatto, Nucl. Phys. **B409** (1993) 257.
11. R. Casalbuoni, A. Deandrea, S. De Curtis, D. Dominici, R. Gatto and M. Grazzini, UGVA-DPT 1995/10-906, hep-ph/9510431.
12. G. Altarelli and R. Barbieri, Phys. Lett. **B253** (1991) 161; G. Altarelli, R. Barbieri and S. Jadach, Nucl. Phys. **B369** (1992) 3; G. Altarelli, R. Barbieri and F. Caravaglios, Nucl. Phys. **B405** (1993) 3.
13. F. Caravaglios, talk at the EPS conference on Particle Physics, Bruxelles, July 1995.
14. R. Casalbuoni, A. Deandrea, S. De Curtis, D. Dominici, R. Gatto and M. Grazzini, Proceedings of the CERN Workshop on Physics at Lep2, CERN Yellow Report, to appear in 1996.